Pregalactic Black Hole Formation with an Atomic Hydrogen Equation of State

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ABSTRACT

The polytropic equation of state of an atomic hydrogen gas is examined for primordial halos with baryonic masses of $M_h \sim 10^7 - 10^9\ M_\odot$. For roughly isothermal collapse around 10^4 K, we find that line trapping of Lyman α (HI and HeII) photons causes the polytropic exponent to stiffen to values significantly above unity. Under the assumptions of zero H_2 abundance and very modest pollution by metals ($< 10^{-4}$ Solar), fragmentation is likely to be inhibited for such an equation of state. We argue on purely thermodynamic grounds that a single black hole of $\sim 0.02 - 0.003 M_h$ can form at the center of a halo for $z = 10 - 20$ when the free-fall time is less than the time needed for a resonantly scattered Lyman α photon to escape from the halo. The absence of H_2 follows naturally from the high, $> 10^4$ K, temperatures that are attained when Lyman α photons are trapped in the dense and massive halos that we consider. An H_2 dissociating UV background is needed if positive feedback effects on H_2 formation from X-rays occur. The black hole to baryon mass fraction is suggestively close to what is required for these intermediate mass black holes, of mass $M_{BH} \sim 10^4 - 10^6\ M_\odot$, to act as seeds for forming the supermassive black holes of mass $\sim 0.001 M_{spheroid}$ found in galaxies today.

Subject headings: cosmology: theory – galaxies: black holes –ISM: clouds – ISM: atoms – atomic processes – radiative transfer

1. Introduction

A fundamental issue in the study of galaxy evolution is the formation of the central (supermassive) black hole. Accretion onto these black holes provides the energy source for active galactic nuclei, which in turn impact the evolution of galaxies (Silk 2005). Earlier attempts at providing seeds for galactic black holes include dynamical friction and collision processes in dense young stellar clusters (Portegies Zwart et al. 2004) and formation from low angular momentum material in primordial disks (Koushiappas, Bullock & Dekel 2004).

In this work, we consider the impact of the polytropic equation of state (EOS) of a metal-free, atomic hydrogen gas on the expected collapse of matter inside massive halos. The impact of a Solar metallicity polytropic EOS on the expected masses of stars in local galaxies has been investigated by Li, Mac Low & Klessen (2003). The influence of molecular hydrogen and metal-poor environments has received detailed interest from, e.g., Abel, Bryan & Norman (2002, 2000); Bromm, Coppi & Larson (2002) for the formation of the first stars and from Scalo & Biswas (2002) and Spaans & Silk (2005) for the properties of the polytropic equation of state. From the work of Li et al. (2003) it has become clear that a polytropic EOS of state, $P \propto \rho^\gamma$ with ρ the mass density and γ the polytropic exponent, strongly suppresses fragmentation of interstellar gas clouds if $\gamma > 1$. This paper concentrates on the impact of Lyman α photon trapping on the EOS and the interested reader is referred to Rees & Ostriker (1977) and Silk (1977) for some of the fundamental thermodynamic and star formation considerations that come into play

here.

2. Model Description

We assume a metal-free hydrogen gas that is cooled by Lyman α emission as it collapses inside a dark matter halo and radiates away about twice its binding energy (Haiman, Spaans & Quataert 2000). Note that Lyman α cooling is expected to dominate over radial contraction factors of at least 15-60 as long as the metallicity is less than 0.1 of Solar (Haiman et al. 2000). The absence of any H_2, that would cool the gas to below 8,000 K, is crucial in this and we come back to this point in Section 4. We further employ a polytropic EOS and a perfect gas law, $P \propto \rho T$, for the gas temperature T, and write γ as

$$\gamma = 1 + \frac{d\log T}{d\log\rho}. \tag{1}$$

This last step is justified (Scalo & Biswas 2002) as long as the heating and cooling terms in the fluid energy equation can adjust to balance each other on a time-scale shorter than the time-scale of the gas dynamics (i.e., local thermal equilibrium). Below, we compute the polytropic EOS for the case that the cooling time is shorter than the free-fall time and for the case that the photon propagation time exceeds the dynamical time.

It should be noted that because γ depends on the (logarithmic) derivative of the temperature with respect to density, it implicitly depends on radiative transfer effects and changes in chemical composition through derivatives of the heating and cooling functions (Spaans & Silk 2000). The Lyman α radiative transfer techniques as described in Haiman & Spaans (1999) and Dijkstra et al. (2006) are used to compute the transfer of Lyman α photons.

We consider spherical dark matter halos that have decoupled from the Hubble flow and are characterized by a mean density of $\rho \approx 200\rho_b(1+z)^3$ at $z = 10 - 20$, for a baryonic number density $\rho_b/m_H = 3 \times 10^{-7}$ cm^{-3} today, hydrogen mass m_H, total halo masses of $M_h = 10^7 - 10^9$ $M_\odot$ and a characteristic size scale of $L = (3M_h/4\pi\rho)^{1/3}$. This yields, over $z = 10 - 20$, a typical mean density and column of $n_0 = 0.05[(1+z)/10]^3$ cm^{-3} and $N_0 = 10^{22}[(1+z)/10]^2(M_h/10^9 M_\odot)^{1/3}$ cm^{-2}, respectively. We further assume that matter inside the halo remains at approximately 10^4 K during its collapse so that an isothermal density profile, $n \propto n_0(L/r)^2$, is applicable for every radius r. Therefore, the column a Lyman α photon has to traverse from a radius r to L scales as $\int_r^L n dr \sim L/r - 1$ with a mass inside of r of $M(r) \sim r$.

In the absence of any ionizing sources, heating is provided by gravitational compression, $\Gamma \propto n^{1.5}$. The velocity dispersion of the gas is thermal and equal to $\Delta V = 12.9 T_4$ km/s, with T_4 in units of 10^4 K. The natural to thermal line width of the Lyman α line is denoted by a and equal to $a = 4.7 \times 10^{-4} T_4^{-1/2}$.

3. Results

3.1. Static Case

With cooling provided by Lyman α emission only, the thermal equilibrium of the baryonic matter in the halo approximately (Spitzer 1978; Haiman et al. 2000) follows, for r' in units of L,

$$7.3 \times 10^{-19} n_e(r') n_H(r') e^{-118,400/T(r')} \epsilon(r') = 1.9 n(r') G M_h / L t_{ff}^{-1}, \tag{2a}$$

with $M_h/L = M_h(r'/L)/r' = M(r')/r'$, electron density n_e, atomic hydrogen density n_H, Lyman α escape fraction ϵ and free-fall time

$$t_{ff} = 4.3 \times 10^7 / n(r')^{1/2} \text{yr}. \tag{2b}$$

It is assumed here that the cooling time at the peak of the cooling curve is $t_c = 3/2nkT_{vir}/n^2\Lambda$, with Boltzmann's constant k, the halo's virial temperature T_{vir} and $\Lambda \approx 2 \times 10^{-22}$ erg s^{-1} cm^3. Depending on the ambient conditions, the medium cools somewhere around the peak of the cooling curve and $T_{vir} < 10^4$ K if the mass is smaller than $10^8[(1+z)/10]^{-1.5}$ $M_\odot$ for the virialization redshift z (Haiman, Rees & Loeb 1997).

The escape fraction ϵ of Lyman α photons from a sphere diminishes from unity when collisional de-excitation above a critical HI column density N_c becomes important (Neufeld 1990). This column N_c depends on the ambient temperature and ionization balance through the probability for collisional de-excitation $p_0 = \frac{q_p n_p + q_e n_e}{A_{21}}$, with proton density n_p, collisional de-excitation rate coefficients q_p and q_e, and the Einstein A coefficient A_{21} connecting the $2p$ and $2s$ states. In this, the

ambient proton density is assumed to be lower than about 10^4 cm^{-3} so that the created $2s$ hydrogen atoms undergo two-quantum decay to the ground state. It follows that N_c ranges between 10^{21} and 10^{23} cm^{-2} for $y = n_p T_4^{-0.17}$ between 10^2 and 10^4 cm^{-3}, respectively (Dijkstra et al. 2006), and is much larger for much smaller proton densities. These values for N_c are a factor of a few larger than the corresponding values for a slab (Neufeld 1990), since resonantly scattered photons escape more easily from a sphere than from a slab for the same surface-to-center optical depth.

Furthermore, following the Monte Carlo radiative transfer techniques in Dijkstra et al. (2006) and Haiman & Spaans (1999), and for a HI column N_H, we find that $\epsilon \approx N_c N_H^{-1.0}$, for $N_H = 2N_c - 100N_c$ and for spherical clouds. In deriving this fit to the numerical results, we have made sure that the line profile is sampled far enough into the wings to accurately determine ϵ and N_c. When applied to a slab, rather than a sphere (see the analytical solution in the appendix of Dijkstra et al. 2006), our method yields results that agree well with those of Neufeld (1990, his Figure 18).

For the resonantly scattered Lyman α line, it follows, for a line center optical depth τ_0, mean line opacity α_s and profile function $\phi(x)$ in normalized frequency units x, that an escaping photon, which scatters N times, experiences a frequency shift $x_s \sim N^{1/2}$ and travels a distance $N^{1/2}/(\alpha_s \phi(x_s))$ that is equal to the size of the medium τ_0/α_s. Hence, $x_s \sim \tau_0 \phi(x_s) \sim (a\tau_0)^{1/3}$ since $\phi \sim a/\pi x^2$. On average, a time $\delta t \sim \frac{L}{c}/(\tau_0 \phi(x_s))$ elapses between the $\sim N$ scatterings. Thus, a time $t_{ph} \sim N\delta t \sim \frac{L}{c}(a\tau_0)^{1/3}$ is required for a photon to escape, where the optical depth is given by $\tau_0 = 1.04 \times 10^{-13} N_H T_4^{-1/2}$. Typically, we have that $\tau_0 > 10^7$. Thus, for a given density, and in the limit that $t_{ff} >> t_{ph}$ with $N_H > N_c$, an increase in column leads to a proportional decrease in the spherical escape probability $1 - e^{-a\tau_0}/a\tau_0 \approx 1/a\tau_0$ and the chance that a scattering hydrogen atom will not suffer collisional de-excitation effects.

Obviously then, Equation (1) has a weak dependence of γ on density for modest columns due to the exponential temperature dependence and the $N_H^{-1.0} \propto n_0^{-2/3}$ scaling of ϵ in the static case and for fixed mass M_h. Under collisional ionization equilibrium, one finds from solving Equations (2) for $T(n)$ that

$$\gamma - 1 \approx -\frac{1}{2log(Cn^{1/2})} = 0.006 - 0.007, \qquad (3)$$

for proton densities larger than 10^{2-4} cm^{-3}, $N_H > 10^{21-23}$ cm^{-2} or $r \leq r_{stat} = 1.0 - 0.01L$ for all halos over $z = 10 - 20$, and where $C \sim 10^{-36}$ $M_h/10^7 M_\odot$ cm$^{3/2}$. Hence, as to be expected, the stiffening of the polytropic EOS is always modest when the exponential temperature dependence of the Lyman α cooling rate acts unchecked.

3.2. Dynamic Case: HI

For the halos considered here, $t_c < t_{ff}$ by a factor of a few. However, if the random walk that a Lyman α photon performs takes a time t_{ph} that is comparable to or longer than the dynamical time on which the halo evolves, cooling is effectively shut down. Photons can then only escape through parts of the line wings that have modest optical depths, while the Lyman α emission becomes zero around line center.

One can show that

$$\epsilon \rightarrow \epsilon \times e^{-\beta t/t_{ff}}, \qquad (4)$$

with $t = t_{ph}$, as more and more photons get trapped in the line core for times exceeding the dynamical time. The multiplier $\beta \sim 2 - 3$ incorporates details of the gravitational collapse of gas shells (e.g. geometry, kinematics) and does not impact our results as long as it does not (or only weakly) depend on density.

That is, the decrease in the number of escaping/cooling Lyman α photons is approximately proportional to the total number of photons somewhere in the line multiplied by the average time a given photon spends in the medium per unit of free-fall time: gravitational collapse scales with $t_{ff} \propto n^{-1/2} \propto r$ and the number of already collapsing shells a photon would have to traverse thus increases linearly in space and time, i.e., $-d\epsilon \sim \epsilon dt/t_{ff}$. It is implicitly assumed here that the scattering–broadened line width, typically larger than 200 km/s (Dijkstra et al. 2006), exceeds any systematic velocity shifts, which is a good approximation for large optical depths. Hence, in Equation (2a) the factor ϵ is now competitive with the

temperature dependence of Lyman α cooling because it picks up an exponential function of density.

Typically, one has $t_{ff} \sim 1.6 \times 10^{15}/n_0^{1/2}$ s and $t_{ph} \sim 5.0 \times 10^{14}/n_0^{1/9}(M_h/10^9 M_\odot)^{1/3}$ s for the adopted halo characteristics. Note in this expression, the weak and negative dependence of t_{ph} on density. This is a consequence of the random walk in both coordinate and frequency space that is performed by the Lyman α photon, yielding a weak $\tau_0^{1/3} \sim n_0^{2/9}$ dependence for t_{ph}, while the size of the halo scales as $n_0^{-1/3}$ for a fixed mass M_h.

The expression for the local thermal balance, Equation (1) above, formally does not change, although all quantities pick up a time dependence, as long as local thermal balance holds. This is still true for $t_{ph} \geq t_{ff}$ and $T \sim 10^4$ K, given that the time needed to thermalize through collisons scales as $1/n$ and the free-fall (heating) time as $1/n^{1/2}$. Similar considerations apply to the ionization balance of hydrogen, but the presence of shocks would require a more careful treatment. The velocity gradients that exist maximally have a magnitude of

$$\delta v \sim r/t_{ff}(r) \sim 10^2 \text{ km/s}, \tag{5}$$

smaller than the scattering broadened line width, and are independent of r if an isothermal density distribution pertains.

One can determine γ straightforwardly for $t \sim t_{ff}$ and a fixed mass M_h. One finds that

$$\gamma - 1 \approx -\frac{\frac{1}{2} + \frac{7}{18} B n^{7/18}}{log(Cn^{1/2}) + Bn^{7/18}}, \tag{6}$$

where it should be noted that $t_{ph}/t_{ff} \propto n^{7/18}$ and that $B \approx 0.5 - 0.1$ cm$^{7/6}$ for $M_h = 10^9 - 10^7$ $M_\odot$ (so $Cn^{1/2} << Bn^{7/18}$).

Evaluation of Equation (6) yields $\gamma - 1 \sim 0.01 - 0.5$ for hydrogen densities of $1 - 10^5$ cm^{-3} for $z = 20$ and a 10^8 $M_\odot$ halo. Note that a density of 1 cm^{-3} is achieved for our halos after a contraction in radius by a factor of a few, much less than the contraction factor $\lambda^{-1} \sim 20$ after which a disk forms (Mo, Mao & White 1998). One finds, for the 10^9 $M_\odot$ halo at $z = 20$ and for the appropriate (column) density scaling with r, e.g., $\tau_0 \sim r^{-1}$ and $n \sim r^{-2}$, that $t_{ff} \geq t_{ph}$ and $\gamma \geq 1.1$ for $r \leq r_{dyn} = 0.02L$. This implies enclosed masses, $M \sim r$, of about $0.02 - 0.003 M_h$ for the adopted isothermal profile and halo masses. The adiabatic value $\gamma = 4/3$ is achieved for $r \leq r_{dyn} = 0.002L$ and a 10^9 $M_\odot$ halo at $z = 20$ (but see see some corrections to γ in the next subsection).

The presence of the 'C term' from Equation (3) does not mean that conversion of Lyman α photons to the two-photon continuum is a significant sink. Rather, for the considered halos, trapping of Lyman α occurs already at densities for which collisional de-excitation by protons is negligible (despite the large columns $N_H \sim N_c$ or somewhat smaller) because the thermal electron abundance is very small[1]. We will come back to the consequences of the rise in temperature associated with $\gamma > 1$ in the next subsection. Finally, the 7/18 dependence on density renders our results relatively insensitive to subtleties in the Lyman α radiative transfer.

3.3. Dynamic Case: Two-Quantum and HeII Corrections

We have assumed that the gas remains close to, but not exactly at, $T = 10^4$ K as far as its density profile is concerned. This is reasonable given the sharpness of the Lyman α cooling function. A value $\gamma > 1$ implies of course that the temperature rises with increasing density, but Lyman α cooling will dominate the local thermal balance for temperatures $T < 5 \times 10^4$ K. Of course, as the temperature rises, so do the electron and proton abundance and this favors the two-photon continuum by decreasing N_c. From Equation (6), thermal ionization balance and the results for $N_c(y)$, one finds that γ weakens towards unity above $\sim 2 \times 10^4$ K when the electron abundance exceeds 0.1. However, this temperature is reached when the density is $10^{3.5}$ cm^{-3} and two-quantum decay is quickly shutdown during the collapse as a density of 10^5 cm^{-3} is exceeded.

In fact, at temperatures above 5×10^4 K HeII line cooling dominates, and the latter also suffers from photon trapping when hydrogen columns exceed 10^{24} cm^{-2}. That is, the HeII Lyman α line at 304Å is similarly opaque (Neufeld 1990), barring the appropriate changes in Einstein A coefficient, elemental abundance etc., as its HI counterpart.

[1]In fact, at densities below 10^3 cm^{-3} one has $N_H << N_c$, but we retain the intuitive form of Equation (6) because the logarithm renders any error insignificant anyway.

One has that $\tau_{He} = 5.2 \times 10^{-14} N(He^+) T_4^{-1/2}$. Hence, photon trapping will continue, for the massive halos that we consider, into the HeII regime at large columns. Also, the much larger HeII Lyman α Einstein A coefficient of $\sim 10^{10}$ s^{-1} boosts the required value of y for a given N_c by two orders of magnitude. The HeII two-photon channel is shut down since densities exceed $10^{5.5}$ cm^{-3} around the HeII cooling peak ($A_{2s1s} \sim 8.2 Z^6$ s^{-1}). In any event, most HeII two-quantum decay photons are absorbed by the (HI Lyman α trapping) neutral hydrogen that surrounds the halo core.

As a result of all this γ remains well above unity and the system evolves adiabatically for densities above $\sim 10^5$ cm^{-3}. Equation (6) provides a good fit to the atomic physics of HI and HeII between $n = 1$ and $n = 10^7$ cm^{-3} if corrected for HI two-photon decay and the change in line opctical depth as HI cooling is superseded by HeII cooling. One finds that

$$\gamma - 1 \approx -\frac{\frac{1}{2} + \frac{7}{18} B'(n) n^{7/18}}{log(Cn^{1/2}) + B'(n) n^{7/18}}, \qquad (7)$$

where $B' = 0$ for $n = 10^3 (M_h/10^9 M_\odot)^{-1/3} - 10^5$ cm^{-3} following Equation (3), and where $B' \approx 0.36B$ if $n \geq n_c = 10^{5.5}$ cm^{-3} and $B' = B$ otherwise. Note here that $t_{ph} > t_{ff}$ always holds and that $B' n^{7/18} >> 1$ for densities larger than n_c, and for all halos, i.e., the HI to HeII switch has a modest impact because the optical depth enters into t_{ph} with a 1/3 power.

4. Discussion and Future Work

The Jeans mass for a 0.1 cm^{-3} halo is about $M_J \sim 3 \times 10^7$ $M_\odot$ at 10^4 K. Over the Lyman α cooling regime, M_J decreases only by a factor of about 8. That is, a $0.02 - 0.003 M_h$ core will likely not experience significant fragmentation during gravitational collapse up to densities of $\sim n(r_c) \approx 10^5$ cm^{-3} at $z = 10 - 20$, after which fragmentation is halted adiabatically. That is, the system cannot cool above a few times 10^5 K either because photons produced by, e.g., bremsstrahlung cannot escape since the bulk of these cooling photons are at energies above a few times 10^{15} Hz and are re-proprocessed into Lyman alpha and trapped in the surrounding neutral exterior of the collapsing cloud. Also, the rise in γ is moderate enough to justify our use of an isothermal density profile.

Although our results are order of magnitude estimates, they connect quite well with the detailed numerical simulations of gravitational collapse by Jappsen et al. (2005) and Klessen, Spaans & Jappsen (2005). The former authors find that a switch to a $\gamma > 1$ region in density space for a collapsing gas sets a characteristic mass scale for fragmentation through the Jeans mass at the ambient density and temperature. Hence, a value $\gamma > 1$ appears to be a robust indicator of the lack of fragmentation. As such, the picture that emerges from detailed hydrodynamical simulations and the shape of the EOS is at least consistent.

Finally, the frequency shift that a Lyman α photon experiences before escape scales approximately as $\nu_{\rm shift} \sim T^{1/6}$ (Dijkstra et al. 2006). Hence, fluctuations in temperature do not strongly influence our results in this respect either.

Thus, allowing for some expected inefficiency, we infer that of order 0.1% of the baryon mass forms a pregalactic black hole of mass $M_{BH} \sim 10^4 - 10^6$ $M_\odot$. Note here that Bromm & Loeb (2003) find a similar inhibition of fragmentation from detailed hydrodynamic simulations that assume a rougly isothermal ($\gamma \sim 1$) collapse. They do not include the trapping effects discussed here. So unless there are numerical resolution effects that play a role, a value of $\gamma = 1$ may already be sufficient to halt fragmentation.

In any case, these so-called intermediate mass black holes (IMBH) are plausible seeds for generating the supermassive, $\sim 0.001 M_{spheroid}$, black holes found in galaxy cores today (c.f. Häring & Rix 2004) by gas accretion (Islam, Taylor & Silk 2003). The inferred presence of pregalactic IMBH and their associated accretion luminosity has been a source of intense speculation with regard to a mechanism for the reionization of the universe (e.g., Madau et al. 2004; Ricotti & Ostriker 2004; Venkatesan, Giroux & Shull 2001). Our results place these speculations on a sounder footing. Moreover, isolated IMBH should exist in galactic halos at a similar mass fraction according to simple models for generating the Magorrian correlation between central black hole mass and spheroid velocity dispersion (Islam, Taylor & Silk 2004; Volonteri, Haardt & Madau 2003) and possibly be detectable as gamma-ray sources (Zhao & Silk 2005).

Still, there are a number of other issues that

should be addressed in the future.

1a) The stiffening of the polytropic equation of state found in this work depends crucially on the absence of any H_2 molecules. For temperatures above 3,000 K this seems plausible because collisional dissociation and charge exchange with H^+ limit the abundance of H_2, while Lyman α trapping keeps the temperature above 10^4 K. Still, H_2 may also form directly in massive halos with virial temperatures above 10^4 K (Oh & Haiman 2002) and reach a universal abundance of $\sim 10^{-3}$. H_2 formation in these cases is a consequence of a freeze-out of the H_2 abundance, in the presence of a large free electron fraction, as gas cools from above 10^4 K on a time scale that is shorter than the H_2 dissociation time. However, in the dense and massive halos that we consider Lyman α trapping, already at modest densities of 1 cm^{-3}, causes the cooling time to increase exponentially, from a level of a few $\times 10^6$ yr at 8,000 K, and to remain larger than the H_2 dissociation time. That is, the gas lingers at 10^4 K, unable to reach 8,000 K or less, and stays at those temperatures because H_2 formation is suppressed (H_2 is easily destroyed by H^+) at the ambient temperatures (Oh & Haiman, reaction 17 on their page 15). As a consequence the H_2 abundance remains at a very low level around 10^4 K, see Figure 4 of Oh & Haiman (2002), and does not contribute to the cooling. Furthermore, Figure 7 of Bromm & Loeb (2003) shows that even a halo with $T_{vir} \sim 10^4$ K (baryonic mass of $\sim 10^7$ $M_\odot$) first heats a large part of the cold (30-100 K) infalling gas to temperatures of $3,000 - 10,000$ K for densities of ~ 1 cm^{-3}, when the formation of, and cooling by, H_2 is incorporated. This is the relevant, minimum temperature range because there is a trough between the Lyman α and H_2 cooling curves here (Oh & Haiman 2002). In this, it is important to realize that our halos are quite massive, baryonic masses of $10^7 - 10^9$ $M_\odot$, and dense ($z > 10$). These values favor Lyman α trapping.

1b) Still, the presence of a UV background from popIII stars, that suppresses the abundance of H_2 molecules (Bromm & Loeb 2003), would certainly be welcome. The critical density for collisional H_2 dissociation to dominate is about 300 cm^{-3}, if a UV background as in, e.g., Bromm & Loeb (2003) is present with which H_2 collisional dissociation has to compete. We do reach this regime early in the collapse so that self-shielding effects would not limit the benefits of H_2 photo-dissociation much (see Bromm & Loeb 2003). Of course, in the absence of a background radiation field, all the time scales (formation, dissociation, etc.) in the system scale as 1/density and thus their ratios are independent of density and the discussion of 1a applies.

1c) The formation of the black hole will introduce a quasar whose power-law spectral energy distribution can boost the formation of H_2 through the H^- route (Haiman, Abel & Rees 2000). Hence, for redshifts below ~ 300, where H^- is no longer destroyed by the CMB, black hole formation as discussed here will facilitate the formation of H_2 and impact the EOS of the gas surrounding the black hole (Scalo & Biswas 2002). The mode that we describe here would then be inhibited unless an H_2 dissociating UV background as in Bromm & Loeb (2003) or Oh & Haiman (2002) is present. Still, the large columns that we consider would shield at least part of the gas from X-ray feedback (see 2b below).

2a) Trace amounts of dust as little as 10^{-4} of Solar are sufficient to absorb all Lyman α photons in a homogeneous halo for the columns considered in this work (Neufeld 1990). Hence, the stiffening of the EOS that we have found disappears once the first metals have been produced because dust emission is optically thin. Also, metals are efficient coolants and, if present, would take over the cooling for radial contraction factors larger than 60 (Haiman & Spaans 1999). Inhomogeneity suppresses dust absorption, but facilitates the escape of Lyman α photons by boosting ϵ (Haiman & Spaans 1999). In any case, the formation of these massive black holes is stopped once the ambient metallicity increases due to star formation. Hence, the fraction of massive primordial galaxies that harbor these black holes is dictated by the fraction of metal-free gas at $z = 10 - 20$. Given that popIII star formation is co-eval with this epoch, the metal-free gas fraction is uncertain and clumpy (Scannapieco, Schneider & Ferrara 2003). Hence, the overall contribution of this mode of black hole formation is somewhat undetermined and likely to lie anywhere between 5 and 50 percent, depending on the proximity of other galaxies. If efficient, this mode may violate the three-year WMAP constraint on the electron scattering optical depth of

$\tau_e \approx 0.09$ (Spergel et al. 2006), because of the large X-ray output expected for these massive and late black holes.

2b) Fortunately, following Ricotti et al. (2005) it is possible to constrain the contribution to τ_e. The latter authors find that τ_e scales as $1/log(N_H)$ and levels off to $\tau_e \sim 0.1$ for columns in excess of 10^{22} cm^{-2}. Given that all our black holes form as the end product of a central collapse from a massive halo, the surrounding gas has column densities between $\sim 10^{22}$ cm^{-2} (from the initial halo masses and redshifts) and $\sim 10^{26}$ cm^{-2} (from the $\gamma = 4/3$ adiabatic points). We have used the models of Meijerink & Spaans (2005) to confirm that these columns are sufficient to re-process X-rays between 1 and 30 keV for metallicities in the surrounding gas between 0 and 10^{-2}. Hence, a value of $\tau_e \sim 0.1$ is appropriate for our black holes, even if they would be the dominant mode of black hole formation.

3) We assume zero angular momentum for the gas, but the radiative transfer is not sensitive to the associated velocity field. Of course, our arguments are purely thermodynamic in nature and do not solve the angular momentum problem if the initial cloud is rotating.

4) We would expect that dwarf spheroidals should have central black holes in the range of $10^4 - 10^6\ M_\odot$, whereas irregulars, and in particular late-forming dwarfs, should not have such central IMBH. In order to substantiate this, more detailed hydrodynamical simulations that include dynamical photon trapping, or its EOS parameterization, should be performed.

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